

Topological (Sliced) Doping of a 3D Peierls System: Predicted Structure of Doped BaBiO₃

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At hole concentrations below $x=0.4$, Ba_{1-x}K_xBiO₃ is non-metallic. At $x = 0$, pure BaBiO₃ is a Peierls insulator. Very dilute holes create bipolaronic point defects in the Peierls order parameter. Here we find that the Rice-Sneddon version of Peierls theory predicts that more concentrated holes should form stacking faults (two-dimensional topological defects, called slices) in the Peierls order parameter. However, the long-range Coulomb interaction, left out of the Rice-Sneddon model, destabilizes slices in favor of point bipolarons at low concentrations, leaving a window near 30% doping where the sliced state is marginally stable.

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I. INTRODUCTION

When a 1D crystal is driven incommensurate, for example, by doping to alter the Fermi wavevector, it is well-known that the modulation of the commensurate crystalline order tends not to be uniform but to accumulate at kinks in the crystalline order parameter [1]. A prototype for this is the Su-Schrieffer-Heeger model [2] for polyacetylene, (CH)_x. In a simple chemical view, the carbon chain has alternating single and double bonds. Crystallography confirms this lattice dimerization. The scalar order parameter (amplitude of the staggered charge disproportionation) has two possible values, $\pm\rho$. Zero-dimensional domain walls (topological defects) separate regions of the 1D chain with positive and negative order parameters. Su *et al.* showed that when carriers are doped into polyacetylene, they create new domain-wall defects and localize into mid-gap electronic soliton states on these defects. Analogous effects in two dimensions [3] occur in layered materials near a Mott insulator phase, such as cuprates, nickelates, and layered manganites. The one-dimensional domain wall in the 2D system is generally called a “stripe.” The 3D analog, which occurs in some cubic manganites [4], is a planar domain wall [5], sometimes called a “sheet” or “lamella.” We prefer the term “slice,” which, like the word “stripe” can be used as verb or noun. Here we discuss a 3D model system in which bonding and strain energies favor slices, but long-range Coulomb interactions prefer point defects. We analyze the system numerically and show that

at doping levels near 30%, Ba_{1-x}K_xBiO₃ is near the crossover where slices may become more favorable than point defects for lowering the energy of the doped-in holes.

Charge-density-wave (CDW) and Peierls systems have scalar order parameters $\rho_Q = N^{-1} \sum_{\ell} \rho_{\ell} \exp(i\vec{Q} \cdot \vec{\ell})$. BaBiO₃ is a simple 3D example. With nominal valence Bi⁴⁺, the Bi 6s-band is half-filled. The material is non-metallic, with a 2eV optical gap [6]. Crystallography [7] shows a doubled unit cell; ρ_{ℓ} alternates with wavevector $Q = (\pi, \pi, \pi)$. The simple cubic sublattice of nominal Bi⁴⁺ ions self-organizes into a bipartite (rocksalt-type) charge-ordered array of nominal Bi³⁺ ($\rho = 4 - \rho_Q$, called “A” sites) and nominal Bi⁵⁺ ($\rho = 4 + \rho_Q$, called “B” sites) ions. The actual value of the order parameter ρ_Q has magnitude ≤ 1 , and takes two degenerate values, $\pm|\rho_Q|$. A simple way to think of this is as a 3D version of a Peierls instability. In a bi-partite lattice with only nearest-neighbor hopping there is accidental Fermi surface nesting: $\epsilon(\vec{k}) = -\epsilon(\vec{k} + (\pi, \pi, \pi))$, and both are zero at the Fermi energy. This guarantees that the crystal can reduce its electronic band energy *via* the electron-phonon interaction by dimerizing. A Hamiltonian which contains this effect was introduced by Rice and Sneddon [8,9].

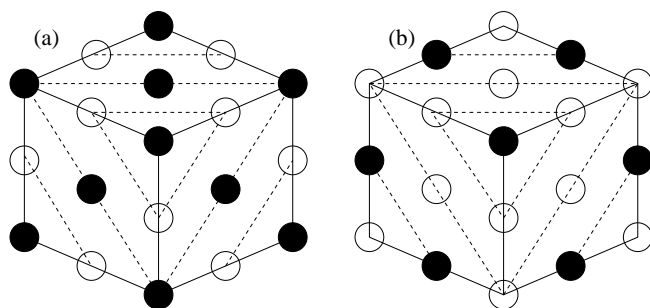


FIG. 1. Fragment of BaBiO₃, showing Bi atoms only. Filled circles are Bi³⁺ (A) ions, and empty circles are Bi⁵⁺ (B) ions. In (a) the normal ABAB... sequence of (111) layers is shown; in (b), a hole-type stacking fault ABBABA..., or slice, is shown.

Regions of charge-ordered BaBiO₃ with $\rho_Q > 0$ are separated from regions with $\rho_Q < 0$ by stacking faults. The simplest stacking fault lies in a (111) plane. In perfectly ordered BaBiO₃, (111) planes are alternating A and B type (Fig. 1a). A stacking fault with no nuclear disorder has either two adjacent A layers (local charge excess -1 per site

on the plane, or electron-doped) or two adjacent B layers (local charge excess 1, or hole-doped) as in Fig. 1b.

Here we point out that the Rice-Sneddon Hamiltonian, given below, predicts that holes or electrons, when introduced, will self-organize into slices. However, the long-range Coulomb interaction, neglected in the Rice-Sneddon model, will destabilize slices except possibly at the most favorable doping level. For light doping, the preferred structure for holes is to self-organize into point defects, small bipolarons, as previously discussed [10–13].

II. RICE-SNEDDON MODEL

The Rice-Sneddon model [8] is simple, well-studied [9,10,14,15], and quite successful [12,13]. In the perovskite crystal structure of BaBiO_3 , each Bi atom is surrounded by six oxygens, and each oxygen is shared by two Bi atoms. At high temperatures (of order 1000K) the crystal is nearly cubic perovskite, but at lower T there are rotations and distortions of the BiO_6 octahedra, which enlarge the unit cell. We believe that the most important effect is the “breathing” displacements $u_0 \approx \pm 0.12\text{\AA}$ [7,17,18] of oxygens away from or toward alternate Bi atoms. These provide a natural electron-phonon mechanism to enhance the Bi 6s charge density on A sites where oxygens breath outward, and reduce the Bi 6s charge density on B sites where the oxygens breathe in. A microscopic Hamiltonian containing the minimal necessary electron-phonon interaction was given by Rice and Sneddon [8],

$$\mathcal{H} = -t \sum_{\langle \ell, \ell' \rangle} c_{\ell}^{\dagger} c_{\ell'} - g \sum_{\ell} e(\vec{\ell}) c_{\ell}^{\dagger} c_{\ell} + \frac{1}{2} K \sum_{\ell, \alpha} u(\vec{\ell} + \frac{\hat{\alpha}}{2})^2. \quad (1)$$

The first term is nearest-neighbor hopping of Bi 6s electrons with hopping integral $t \approx 0.35$ eV. The index of summation ℓ implicitly includes a spin as well as site quantum number. The filling is $1 - x$ electrons per site. The variable $u(\vec{\ell} + \hat{\alpha}/2)$ (with $\alpha = x, y, z$) is the displacement along a Bi-O-Bi bond in the $\hat{\alpha}$ direction of the oxygen located at position $(\vec{\ell} + \hat{\alpha}/2)a$. The variable $e(\ell)$ is the local dilation or “breathing” amplitude of the 6 oxygens which surround the Bi ion at site ℓ ,

$$e(\ell) = [u(\vec{\ell}, x + \hat{x}/2) - u(\vec{\ell} - \hat{x}/2)] + [x \rightarrow y] + [x \rightarrow z]. \quad (2)$$

The Einstein restoring force $K \approx 19$ eV/ \AA^2 is fitted to the measured 70 meV frequency of the Raman-active Peierls breathing mode [19].

At half-filling, this model opens a Peierls gap at the Fermi level; the electron-phonon interaction parameter $g \approx 1.39$ eV/ \AA is fitted to the measured [6] ≈ 2 eV gap. This is an ordinary size electron-phonon coupling. The change in Coulomb field of a charge $-2e$ oxygen ion gives $g \approx 1.2$ eV/ \AA . The resulting dimensionless coupling constant $\Gamma \equiv g^2/Kt$ is ≈ 0.30 , intermediate between the weak ($\Gamma < 0.2$) and strong ($\Gamma > 0.4$) coupling regimes. In this middle regime, neither hopping nor electron-phonon energy dominates [12]. The ground-state of undoped BaBiO_3 is as close to a bipolaronic crystal (large Γ , $|\rho_Q| \approx 1$) as to the conventional Peierls-CDW (small Γ). We calculated [12] the order parameter ρ_Q at $\Gamma = 0.3$ to be 0.82. The corresponding oxygen displacement $u_0 = 2g\rho_Q/K = 0.12\text{\AA}$ agrees with the diffraction and EXAFS measurements [7,17,18], showing that the model is internally consistent.

We previously found that excitations across the Peierls gap form self-trapped excitons [13]. We also reported that holes inserted into BaBiO_3 self-trap and form bipolarons [12] since the coupling strength exceeds $\Gamma_c = 0.17$. These are doubly charged point defects, corresponding to local depressions of the order parameter where the oxygen distortion $e(\vec{\ell}) \rightarrow 0$ for $t = 0$. For non-zero hopping t the bipolaron spreads out and evolves continuously from a small bipolaron ($\Gamma \gg \Gamma_c$) to large CDW-like bipolaron as $\Gamma \rightarrow \Gamma_c$. The stability of bipolaron defects provides a simple explanation why dilutely doped BaBiO_3 remains insulating and diamagnetic.

The “disproportionation reaction” $2\text{Bi}^{4+} \rightarrow \text{Bi}^{3+} + \text{Bi}^{5+}$ has been much discussed in the literature on BaBiO_3 . Using the definition $2U = E(\text{Bi}^{3+}) + E(\text{Bi}^{5+}) - 2E(\text{Bi}^{4+})$, one can say that the effective Hubbard U parameter is negative. Two factors may contribute to U , electron-phonon effects (expected to be attractive) and Coulomb repulsion ($U = U_{\text{ep}} + U_{\text{el}}$). If one wants to assign the mechanism for disproportionation completely to Coulomb interactions, then the Hubbard U calculated with all atoms held stationary in cubic perovskite positions (defined as U_{el}) should be negative. Vielsack and Weber [16] did careful calculations of U_{el} , finding no evidence for negative values, but instead a small positive value $U_{\text{el}} \approx 0.6 \pm 0.4$ eV. Therefore we shall temporarily ignore the on-site Coulomb repulsion U_{el} . Apparently the Bi^{4+} ion must reorganize its environment in order to stabilize the charge disproportionation.

III. HOLE DOPING

What happens at finite doping concentrations x ? Assuming sufficient electron-phonon coupling

to destroy the undistorted metal, there are two possibilities, (1) bipolarons, and (2) slices.

(1) Numerical studies by Yu *et al.* [10], and confirmed by us, show that randomly located point bipolarons are at least metastable. Depending on whether bipolarons attract or repel, the system could then either phase-separate into undoped and doped regions, or form spatially separated bipolarons. When bipolarons are small, the energy of an array of bipolarons is described by a pairwise additive potential $V(\Gamma, r)$, containing the repulsive long-range Coulomb interaction V_{Coul} (neglected for the time being) and the interaction V_t , induced by hopping. V_t decays exponentially with the distance r between two bipolarons (like bipolaron wave functions.) Since bipolarons can only sit on former A sites, the nearest neighbor interaction $V_0 = V_t(\Gamma, \sqrt{2}a)$, which is also the strongest interaction, is between bipolarons separated by $a(110)$. We computed V_0 numerically by optimizing the oxygen-positions self-consistently for given bipolaron positions. Fig. 2 shows that large CDW-like bipolarons attract each other, whereas small bipolarons repel, with $V_t \rightarrow 0$ in the atomic limit $t \rightarrow 0$. Perturbation theory around $t = 0$ gives a bipolaron repulsion. At the physically relevant value $\Gamma = 0.3$, bipolarons repel according to Fig. 2, but multi-bipolaron interactions become important with decreasing Γ . Therefore, in the intermediate coupling regime we must rely on exact numerical diagonalization of Eq. (1) [10].

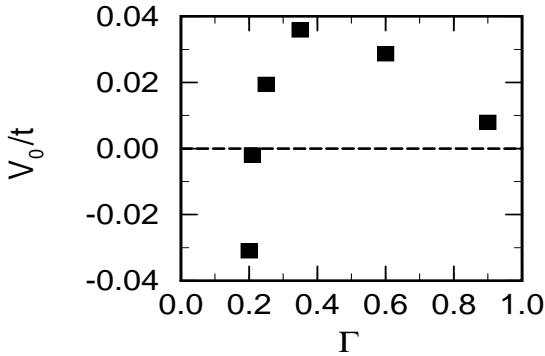


FIG. 2. Interaction potential V_0 for two bipolarons sitting on $(0, 0, 0)$ and $(a/2)(1, 1, 0)$. Large bipolarons attract whereas small bipolarons repel.

(2) In contrast to bipolaron defects where the order parameter never changes sign, holes could form topological defects. Phase-slips are planar defects with sign-changes of the charge order parameter ρ_Q and of the breathing order parameter $\hat{e} = (-1)^{t_{el}}$. Consider a “BB” stacking fault in the 111-direction (Fig. 1b). In the atomic ($t \rightarrow 0$) limit, each B-site on a phase-slip plane has three displaced and three undisplaced oxygen neighbors, i.e. $e_l \approx \pm 3u_0$. Phase-slips accumulate one hole for every two atoms on a (111) BB bilayer. The

average hole charge on phase-slip B-sites is thus $\rho_{\text{hole}} \approx +\frac{1}{2}$. The actual values of hole charge found for $\Gamma = 0.3$ and $x = 1/4$ are shown in Fig. 3.

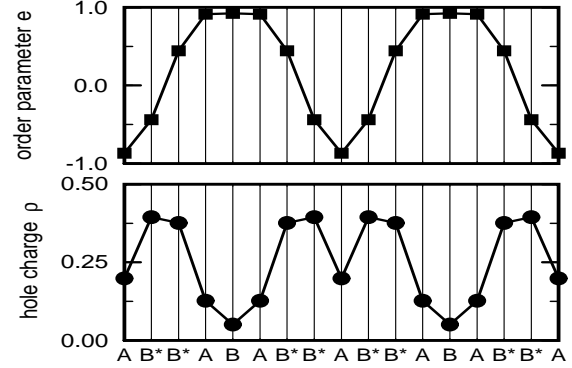


FIG. 3. Normalized breathing order parameter $e/6u_0$ (upper panel) and hole charge ρ_{hole} (lower panel) for a $x = 1/4$ -doped phase-slip structure $(AB^*B^*ABAB^*B^*)_N$ at $\Gamma = 0.3$ (two periods shown). The order parameter changes sign across each phase-slip B^*B^* . Holes are localized on phase-slip planes with slight spreading to neighboring A-planes. The hole charge is computed relative to a reference system with the same structure, $(AB^*B^*ABAB^*B^*)_N$, but with A charges fixed at their undoped value $4+\rho_Q$, B charges fixed at $4-\rho_Q$, and B^* charges fixed at $4-\rho_Q/2$ corresponding to zero doping, with the charge deficit ρ_Q spread equally on the two adjacent B^* layers. The charge fluctuation ρ_Q has the value 0.82 at $\Gamma = 0.3$.

Our aim is to find the ground-state hole arrangement, testing the stability of several bipolaronic (1) versus the phase-slip (2) solutions numerically. We did a series of calculations on $x = 1/4$ -doped systems. The bipolaron systems (1) were (a) maximal spacing between bipolarons, obtained when they occupy center and corner sites of a tetragonal 16 atom unit cell (*bct* structure); (b) a simple cubic (*sc*) arrangement of bipolarons sitting at the corner sites of a 8 atom cubic cell; (c) a disordered structure with random bipolaron positions; (d) phase-separated structures based on unit cells containing 8 or 16 (111) planes, where we replace one or two near neighbor A planes with bipolaronic B planes, $[(AB)_3B_2]_N$ or $[(AB)_6B_4]_{N/2}$. We also looked at unit cells with 8 or 16 111 planes containing phase slips (2) $[AB(ABB)_2]_N$ and $[(AB)_2(ABB)_4]_{N/2}$. Finally, we looked at the undistorted metal $[C]_N$ where each atom C has a nominal $\text{Bi}^{4.25+}$ valence.

Phase separation was strongly disfavored, while separated bipolarons and phase slips were all metastable. The phase slips weakly repelled, preferring the 8 plane solution to the 16 plane solution. The order parameter and hole charge density of this 8 plane solution are shown in Fig.(3). The stability is determined computing the total energy $E_{\text{tot}}(\{u_i\})$ given by Eq.(1), which is a function of the oxygen-positions u_i . We start by guessing oxy-

gen positions to get close to a local minimum in the energy landscape, then vary oxygen positions using a gradient minimization routine to find a self-consistent minimum. For smaller periodic structure, we used k -space sampling in the corresponding Brillouin zones (8000 k -points). For each \vec{k} , \mathcal{H} is diagonalized exactly. States are filled with two electrons up to the desired doping. For the random bipolaron structure we used large asymmetric clusters (≈ 400 atoms) with periodic boundary conditions and $k = 0$ only. Initial oxygen-positions had Peierls order with small random deviations. Our calculations on random configurations generally reproduce the earlier calculations of Yu *et al.* [10].

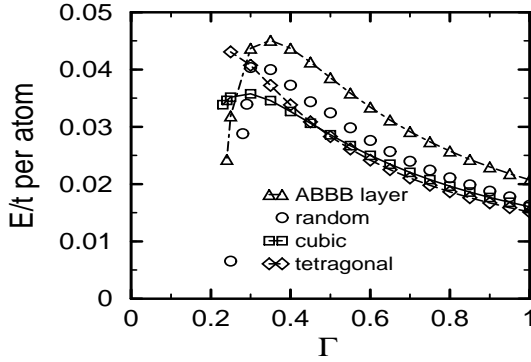


FIG. 4. Energy difference E/t of bipolaron structures (1a-d) relative to the most favorable phase-slip structure (2) for $x = 1/4$ -doping.

Energies of various states at doping $x=1/4$ are shown in Fig. 4. In the atomic limit $t = 0$ all bipolaron and phase-slip structures are degenerate. All that is required is that no A site should have an A first neighbor. Below a critical coupling strength $\Gamma_c(x)$ (≈ 0.2 for doping $x=1/4$), distorted structures become unstable with respect to the undistorted metallic structure. Above $\Gamma_c(x)$, we find numerically that at $x = 1/4$ holes strongly prefer to order in phase-slips (2). Bipolaronic structures (1) are quite similar to each other in energy, and behave as expected from the bipolaron-pair interaction. In the intermediate coupling range there occurs a cross-over from the tetragonal to the layered BBB-structure as stable bipolaron configuration, corresponding to a change in the overall bipolaron interaction from weakly repulsive to weakly attractive. For $1/4$ -doped BaBiO_3 , the energy gain for the phase-slip solution compared to the bipolarons is about 50-65 meV per hole.

Thus, contrary to previous studies of doping of this model [9,10], we find that the stable doping state is not bipolarons but phase slips. Does this model correspond sufficiently to reality for BaBiO_3 ? Iwano and Nasu [14] use a more complicated hopping and electron-phonon interaction,

but we believe that such corrections are not the relevant ones. There are also (i) small structural distortions (rotations of oxygen octahedra) beyond the breathing-mode distortions considered here, and (ii) non-adiabatic effects (such as zero point motion) associated with the fact that the oxygen mass is not infinitely large compared with the electron mass. We believe that both of these also have little relevance. It is harder to dismiss two other effects: (iii) the disorder caused by the dopant atoms, and (iv) the long-range Coulomb interaction, both omitted so far. Of these, the last is clearly important, as we now show, and tends to destabilize phase-slips relative to distributed bipolarons.

IV. LONG-RANGE COULOMB EFFECT

At low doping, there is a large Coulomb cost in putting charges onto stacking faults instead of widely distributed point charges. We modeled this as follows. The Madelung energy was computed by the Ewald method for $(\text{Ba}^{2+})_2(\text{Bi}^{3+}\text{Bi}^{5+})(\text{O}^{2-})_6$. The calculation was re-done for many large unit cells, with holes added on selected Bi ions ($\text{Bi}^{3+} \rightarrow \text{Bi}^{5+}$), and compensating negative charges distributed uniformly throughout space. This is one way to mimic potassium doping of BaBiO_3 . First consider the sliced state. By numerical calculation for uniformly distributed phase slips at many values of x between $1/39$ and $1/3$, we found a good fit to the formula $E_S = (-25.1 + \pi/9x)(e^2/2a\epsilon_\infty)$. The static electronic screening $\epsilon_\infty \approx 5$ was measured by Tajima *et al.* [20]. E_S is the difference of energy per hole between the sliced solution and the undoped Peierls insulator. The term $\pi/9x$ is the analytic result for idealized uniform sheets of charge of vanishing thickness, arranged periodically in a compensating charge background. The term -25.1 corrects for the discreteness of the charges, the absence of the self-interaction, and includes the Coulomb energy of the holes with the background BaBiO_3 lattice.

We also need the energy difference per hole E_B of optimally spaced bipolarons relative to the undoped Peierls insulator. For doping $x = 1/n^3$ the bipolarons can be placed on a sublattice of face-centered cubic (*fcc*) form, with maximum spacing. Numerical results for $n=2, 3$, and 4 fitted well to the formula $E_B = (-21.2 - 4.585x^{1/3})(e^2/2a\epsilon_\infty)$. The term $-4.585x^{1/3}$ is the Madelung energy of an *fcc* lattice of charges $2e$ in a uniform compensating background, and the constant -21.2 accounts for the energy of the holes in the background BaBiO_3 lattice. The difference $E_S - E_B$ is the Coulomb penalty per hole for forming slices or phase slips. It is plotted in Fig 5, becoming large at low x with

a shallow minimum near $x = 0.3$.

A more accurate estimate of the Coulomb penalty is shown as an **x** on Fig 5, and was obtained by doing Ewald sums using the actual Bi site charges ($\rho_i = 5 - n_i$ where $n_i = 2 < \psi_i^\dagger \psi_i >$ is the local occupancy of the Bi s states) and the actual $(AB^*B^*ABAB^*B^*)_N$ slice structure and bct bipolaron structure. It is seen on the figure that at 1/4 doping, the Coulomb penalty is twice as high as the Peierls benefit in forming slices. These calculations did not account for the actual random positions of the compensating negative charges where Ba^{2+} ions are replaced by K^+ ions. Instead, the compensating charge was distributed uniformly. The calculations were repeated for a model where the compensating negative charges were equally shared by each Ba atom ($Ba^{2+} \rightarrow Ba^{(2-x)+}$). This only changed the constant terms in E_S and E_B , but did not affect the difference, $E_S - E_B$ plotted here.

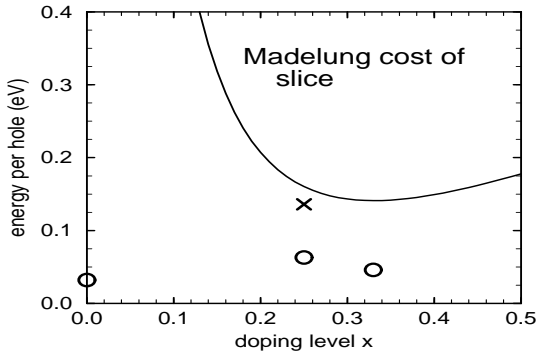


FIG. 5. The curve is an analytic approximation to the Madelung or Coulomb repulsion of a periodic array of (111) phase-slips (slices) compared with optimally spaced separated bipolarons at the same average density, where doped-in holes are modelled as point charges (charge $=e$) with screening ϵ_∞ set to 5. The open circles give the numerically computed Peierls short-range stabilization energy (from hopping and electron-phonon energy) of sliced solutions relative to separated bipolarons. The symbol **x** is the exact numerical Madelung energy difference at 1/4 doping between the sliced solution and the bct bipolaron solution with the numerically computed layer charges ρ_i .

The Madelung sums discussed above are not the complete Coulomb effect. The missing on-site repulsion has the opposite effect, preferring sliced solutions to distributed bipolarons. If we add back the local term

$$\mathcal{H}_U = U_{el} \sum_i (\rho_i - 4)^2 \quad (3)$$

and treat it as a first-order perturbation, the sliced solution at 1/4-doping has lower on-site energy by $0.121U_{el}$ per hole than the bct bipolaron solution. Using the value $U_{el} \approx 0.6 \pm 0.4$ eV from Vielsack and

Weber [16], the onsite correction 0.07 ± 0.05 eV is potentially sufficient to re-stabilize the sliced solution. Thus we can say that the best place to look for sliced structures in $BaBiO_3$ is in the range near 1/4 to 1/3 doping, but we cannot predict whether the sliced solution will be destroyed by disorder or Coulomb effects.

Doping 1/3 is the highest at which a simple sliced solution $((ABB)_N)$ is possible. For higher doping, either the A-planes acquire polaron defects or the sliced solution is destroyed and purely bipolaronic states win.

It is worth mentioning that there have been reports [21] (subsequently attributed to electron beam heating effects [22]) of superlattice diffraction spots in doped $BaBiO_3$, not apparently identical to the superstructures predicted here. A further search in the doping region $0.2 < x < 0.35$ would be interesting.

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